

# Duality Symmetries and Supersymmetry Breaking in String Compactifications\*

**Gabriel Lopes Cardoso and Dieter Lüst**

*Humboldt-Universität zu Berlin  
Institut für Physik  
D-10099 Berlin<sup>†</sup>*

and

**Thomas Mohaupt**

*DESY-IfH Zeuthen  
Platanenallee 6  
D-15738 Zeuthen<sup>‡</sup>*

## ABSTRACT

We discuss the spontaneous supersymmetry breaking within the low-energy effective supergravity action of four-dimensional superstrings. In particular, we emphasize the non-universality of the soft supersymmetry breaking parameters, the  $\mu$ -problem and the duality symmetries.

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<sup>†</sup>e-mail addresses: GCARDOSO@QFT2.PHYSIK.HU-BERLIN.DE, LUEST@QFT1.PHYSIK.HU-BERLIN.DE

<sup>‡</sup>e-mail address: MOHAUPT@HADES.IFH.DE



# 1 Introduction

Based on theoretical motivations, in particular the so-called hierarchy problem, and stimulated by some indirect experimental hints, like coupling constant unification and the top quark mass, the minimal supersymmetric standard model (MSSM) was extensively discussed during the last years. Unfortunately, the necessary violation of supersymmetry has to be put in by hand into the MSSM and is described by the so-called soft supersymmetry breaking parameters (SSBP) like the gaugino masses etc. For reasons of simplicity these SSBP were assumed in most of the phenomenological discussions to be universal for all different gauginos and also for the various matter fields. For some SSBP, a possible deviation from universality is severely constrained by phenomenological requirements like the absence of flavor changing neutral currents [1].

On the other hand, superstring theories are a very promising candidate for a consistent quantization of gravity. For this purpose, the typical string scale has to be identified with the Planck mass  $M_P$  of order  $10^{19}$  GeV. Therefore one strongly hopes that superstrings may solve some puzzles concerning quantum physics at  $M_P$ . Now for the actual relevance of superstring theories it is of most vital importance to make direct contact to the standard model (SM) or perhaps better to the MSSM. This program attracted a lot of attention during the last 10 years, and the results of this research are, at least conceptually, quite successful. Indeed, the low energy effective lagrangian of a large class of four-dimensional heterotic string theories is just given by the standard  $N = 1$  supergravity action with gauge group potentially containing the gauge group of the SM and with matter coming very near to the three chiral families of the SM. Deriving the effective string action, it is very important to realize that the low energy spectrum and the low energy effective interactions among the almost massless fields are to some extent controlled by the stringy symmetries which are still reminiscent after integrating out the infinite number of massive modes. A particular nice example of this kind are the well established duality symmetries (for a review see [2]) which proved to provide useful information about the effective string action on general grounds.

A very attractive feature of  $N = 1$  supergravity in general is the fact that upon spontaneous supersymmetry breaking in some hidden sector of the theory the SSBP in the observable sector automatically emerge due to gravitational couplings among observable and hidden fields. Thus in string theory the SSBP are, at least in principle, calculable from first principles. However at the moment, the actual mechanism of supersymmetry breaking is far from being completely understood. However recently it was demonstrated [3, 4, 5, 6, 7] that, parametrizing the SSBP without specifying the actual supersymmetry breaking mechanism, some interesting generic features of supersymmetry breaking in superstring theories can be derived. In particular it turned out that the SSBP are generically non-universal [3].

This contribution will be organized as follows: first we will set up the general formalism of supersymmetry breaking in  $N = 1$  supergravity with special emphasis on the structure of the SSBP in four-dimensional strings. As more specific examples we will

then present some results for Abelian orbifolds.

## 2 $N = 1$ effective supergravity action for four-dimensional heterotic strings

Let us first specify the string modes with masses small compared to  $M_P$  which we assume to appear in the effective action. First there is the  $N = 1$  supergravity multiplet containing the graviton field and the spin  $\frac{3}{2}$  gravitino. Next, the gauge degrees of freedom are described by  $N = 1$  vector multiplets  $V_a$  with spin 1 gauge bosons and Spin  $\frac{1}{2}$  gauginos  $\lambda_a$ . The gauge index  $a$  is assumed to range over the SM gauge group  $SU(3) \times SU(2) \times U(1)$  and an unspecified hidden gauge group  $G_{\text{hid}}$ . Finally we consider chiral matter multiplets  $\Phi^I$  with complex scalars and spin  $\frac{1}{2}$  Weyl fermions. These chiral fields, i.e. the index  $I$ , separate into so-called matter fields  $Q^\alpha$  which contain the matter of the MSSM,  $Q_{\text{SM}}^\alpha = (q, l, H_1, H_2)$ , and matter which only transforms non-trivially under  $G_{\text{hid}}$ ,  $Q_{\text{hid}}^\alpha$ . The second type of chiral fields  $\Phi^I$  correspond to the so-called moduli fields  $M^i$  whose vacuum expectation values (vev's) are undetermined in perturbation theory since the  $M^i$  correspond to the free parameters of the four-dimensional string models. The moduli are assumed to be SM singlets (however note that  $H_1$  and  $H_2$  could be in principle moduli). The duality group  $\Gamma$  acts on the moduli  $M^i$  as discrete reparametrizations,  $M^i \rightarrow \tilde{M}^i(M^i)$ , which leave the underlying four-dimensional string theory invariant. Therefore, the effective action of the massless field must be  $\Gamma$  invariant which provides a link between  $L_{\text{eff}}$  and the theory of  $\Gamma$ -modular functions [8]. Moreover strong restrictions on the massless spectrum arise [3] due to the required absence of potential duality anomalies.

The effective  $N = 1$  supergravity action, up to two space-time derivatives, is specified by three different functions of the chiral fields  $\Phi^I$  [9]. First, the Kähler potential  $K$  is a gauge-invariant real analytic function of the chiral superfields. To compute later on the SSBP it is enough to expand  $K$  up to quadratic order in the matter fields:

$$K = K_0(M, \bar{M}) + K_{\alpha\bar{\beta}}(M, \bar{M})Q^\alpha\bar{Q}^{\bar{\beta}} + \left(\frac{1}{2}H_{\alpha\beta}(M, \bar{M})Q^\alpha Q^\beta + \text{h.c.}\right) \quad (2.1)$$

Note that for SM matter fields the last term in eq.(1) can be non-vanishing only for a mixing term of the two Higgs fields:  $(\frac{1}{2}H_{12}(M, \bar{M})H_1H_2 + \text{h.c.})$ .  $K_0$  is just the Kähler potential of the Kählerian moduli space  $\mathcal{K}_0$ .  $\Gamma$ -duality transformations act as Kähler transformations on  $K$ ,  $K \rightarrow K + g(M) + \bar{g}(\bar{M})$  ( $g(M)$  is a holomorphic function of the moduli), and induce a ‘rotation’ on the matter fields,  $Q^\alpha \rightarrow h_{\alpha\beta}(M)Q^\beta$ .

Next we consider the moduli dependent effective gauge couplings  $g_a(M, \bar{M})$ :

$$\begin{aligned} g_a^{-2}(M, \bar{M}) = \text{Re } f_a(M) - \frac{1}{16\pi^2} & \left( (C(G_a) - \sum_\alpha T_a(\alpha)) K_0(M, \bar{M}) \right. \\ & \left. + 2 \sum_\alpha T_a(\alpha) \log \det K_{\alpha\bar{\beta}}(M, \bar{M}) \right). \end{aligned} \quad (2.2)$$

$C(G_a)$  is the quadratic Casimir of the gauge group  $G_a$  and  $T_a(\alpha)$  the index of the massless matter representations. The holomorphic gauge kinetic function  $f_a(M)$  includes the tree level moduli dependence as well as possible one-loop quantum corrections from massive modes; however beyond one loop there are no perturbative corrections to  $f_a(M)$  [10]. The non-holomorphic terms in eq.(2) originate from one-loop corrections involving massless fields. Specifically, these terms describe the presence of Kähler as well as  $\sigma$ -model anomalies [11, 12].  $g_a(M, \bar{M})$  has to be a duality invariant function. Therefore the duality non-invariance of the non-holomorphic anomaly terms has to be cancelled by a non-trivial transformation behaviour of  $f_a(M)$ :

$$f_a(M) \rightarrow f_a(M) + \frac{1}{8\pi^2}((C(G_a) - \sum T_a(\alpha))g(M) - \frac{1}{4\pi^2} \sum T_a(\alpha) \log \det h_{\alpha\beta}(M)). \quad (2.3)$$

Third the superpotential will be conveniently split into a SUSY-preserving tree-level part and into a SUSY-breaking piece which does not depend on the matter fields:

$$W = W_{\text{tree}}(Q, M) + W_{\text{SUSY-breaking}}(M). \quad (2.4)$$

Duality invariance of the effective action demands that  $W$  transforms as  $W \rightarrow e^{-g(M)}W$ . The structure of  $W_{\text{tree}}$  is such that it generates the moduli-dependent Yukawa couplings for the matter fields as well as possible moduli-dependent mass terms for some hidden matter fields; the observed matter fields are assumed to stay massless for all values of the moduli fields:

$$W_{\text{tree}} = \frac{1}{3}h_{\alpha\beta\gamma}(M^i)Q^\alpha Q^\beta Q^\gamma + \frac{1}{3}h_{i\alpha\beta}(M^i)Q_{\text{hid}}^\alpha Q_{\text{hid}}^\beta. \quad (2.5)$$

Thus it may happen that at some points in the moduli space,  $h_{i\alpha\beta}(M^i) = 0$ , there are additional massless hidden matter fields. Very often they go together with additional massless gauge bosons at these points.

Essentially, there are two very promising mechanisms of supersymmetry breaking in the last years' literature. First at tree level by the so-called Scherk-Schwarz mechanism [13]. This can be described in the effective field theory by a tree-level superpotential. Second supersymmetry can be broken due to non-perturbative effects. Unfortunately it is not possible at the moment to calculate these non-perturbative effects directly in string theory. However, let us assume that non-perturbative field theory effects give a dominant contribution to the spontaneous breaking of supersymmetry. In particular, one can show that non-perturbative gaugino condensation in the hidden gauge sector potentially breaks supersymmetry [14]. Integrating out the dynamical degrees of freedom corresponding to the gaugino bound states, the duality invariant [15, 16] gaugino condensation can be described by an effective non-perturbative superpotential, which depends holomorphically on the moduli fields:

$$W_{\text{SUSY-breaking}}(M) = e^{\frac{24\pi^2}{b_a} f_a(M)}, \quad (2.6)$$

( $b_a$  is the  $N = 1$   $\beta$ -function coefficient). It is remarkable that this expression is in a sense exact since  $f_a(M)$  is only renormalized up to one loop. It is this exactness of

$W_{\text{SUSY-breaking}}$  which provides very strong confidence in the applicability of the used method.

Now let us discuss the form of the SSBP in the effective action which arise after the spontaneous breaking of local supersymmetry. This discussion will not refer to the actual (perturbative or non-perturbative) breaking mechanism; nevertheless some interesting information about these couplings can be obtained at the end. The scalar potential in the low-energy supergravity action has the form [9]

$$V = |W_{\text{SUSY-breaking}}(M)|^2 e^{K_0} (G^i G_i - 3). \quad (2.7)$$

( $e^G = |W|^2 e^K$ ,  $G_I = \frac{\partial G}{\partial \phi_I}$ .) Deriving this formula we have assumed that, upon minimization of  $V$ ,  $\langle G_\alpha \rangle = 0$  and  $\langle Q_\alpha \rangle = 0$  in the matter sector. This assumption, which is satisfied in most realistic scenarios, means that the spontaneous supersymmetry breaking takes places in the moduli sector, i.e.  $\langle G_i \rangle \neq 0$  for at least one of the moduli fields. Then the gravitino mass becomes

$$m_{3/2} = e^{K_0(M, \bar{M})/2} |W_{\text{SUSY-breaking}}(M)|. \quad (2.8)$$

$m_{3/2}$  should be of order TeV; thus the smallness of this scale compared to  $M_P$  must come either from the Kähler potential and/or from the superpotential. Now we obtain the following SSBP: first the gaugino masses take the form

$$m_a(M, \bar{M}) = \frac{1}{2} m_{3/2} G^i(M, \bar{M}) \partial_i \log g_a^{-2}(M, \bar{M}). \quad (2.9)$$

The scalar masses (squarks and sleptons) become [4]

$$m_{\alpha\bar{\beta}}^2 = m_{3/2}^2 [K_{\alpha\bar{\beta}}(M, \bar{M}) - G^i(M, \bar{M}) G^{\bar{j}}(M, \bar{M}) R_{i\bar{j}\alpha\bar{\beta}}]. \quad (2.10)$$

( $R_{i\bar{j}\alpha\bar{\beta}} = \partial_i \bar{\partial}_{\bar{j}} K_{\alpha\bar{\beta}} - \Gamma_{i\alpha}^\gamma K_{\gamma\bar{\delta}} \bar{\Gamma}_{\bar{j}\bar{\beta}}^{\bar{\delta}} \Gamma_{i\alpha}^\gamma = K^{\gamma\bar{\delta}} \partial_i K_{\alpha\bar{\gamma}}$ . These parameters are generically of the order of  $m_{3/2}$ . Their exact values depend on the details of  $K$ ,  $W$  and the (dynamically fixed) vev's of the moduli fields. It is quite evident that in general these SSBP are non-universal [3]. The non-universality arises due to the non-universal moduli dependence of the gauge couplings and the matter kinetic energies. Similar expression can be also obtained for the trilinear couplings [17, 4, 5]

Finally let us investigate the possible appearance of a mass mixing term for the two standard model Higgs fields  $H_1$ ,  $H_2$  which is necessary for the correct radiative breaking of the electro-weak gauge symmetry. Clearly, a tree-level mixing due to a quadratic term in the superpotential,  $W_{\text{tree}} = \mu H_1 H_2$ , would be a disaster, since it will be most likely of the order of  $M_P$ . (This is often called the  $\mu$ -problem.) However, if there exist [18] a possible, holomorphic mixing term  $H_{12}$  among  $H_1$  and  $H_2$  in the tree-level Kähler potential (see eq.(1)), then an effective  $\mu$ -term will be generated after the spontaneous breaking of supersymmetry:

$$W_{\text{eff}} = \hat{\mu} H_1 H_2, \quad \hat{\mu} = m_{3/2} [H_{12}(M, \bar{M}) - G^{\bar{i}} \partial_{\bar{i}} H_{12}(M, \bar{M})]. \quad (2.11)$$

### 3 Abelian orbifolds

In this chapter we want to apply our previous formulas to the case of Abelian orbifold compactifications [19]. Every orbifold of this type has three complex ‘planes’, and each orbifold twist  $\vec{\theta} = \theta_i$  ( $i=1,2,3$ ) acts either simultaneously on two or all three planes. Generically, for all four-dimensional strings there exist as moduli fields the dilaton ( $D$ ) – axion ( $a$ ) chiral multiplet  $S = e^D + ia$ . Then the tree-level Kähler potential for the  $S$ -field has the form  $K_0 = -\log(S + \bar{S})$ .

Next we consider the internal moduli of the orbifold compactification. We will concentrate on the untwisted moduli fields. For each Abelian orbifold there exist at least three Kähler class moduli  $T_i$  each associated to one of the three complex planes. We will call the  $T_i$  (2,2) moduli, since they do not destroy a possible (2,2) superconformal structure of the underlying string theory, i.e. their vev’s do not break the (2,2) gauge group  $E_6 \times E_8$ . Next we consider so-called (0,2) untwisted moduli which are generically present in any orbifold compactification. A non-vanishing vev for these kind of fields destroys the (2,2) world sheet supersymmetry and breaks  $E_6 \times E_8$  to some non-Abelian subgroup. In addition they will generically give mass to some matter fields by a superpotential coupling. Specifically these types of moduli correspond to continuous Wilson line background fields [20] which are again associated to each of the three complex planes. For the case that  $\theta_i \neq \pm 1$ , there is generically at least one complex Wilson line field  $A_i$  (for example a  $\underline{27}$  of  $E_6$ ). The combined  $T_i, A_i$  Kähler potential reads [21, 22]

$$K_0 = -\log(T_i + \bar{T}_i - A_i \bar{A}_i) \quad (3.12)$$

and leads to the Kähler metric of the space  $\mathcal{K}_0 = SU(1, 2)/SU(2) \times U(1)$ . If  $\theta_i = \pm 1$  there will be additional moduli fields namely, first, the (2,2) modulus  $U_i$  which corresponds to the possible deformations of the complex structure. In addition there will be again some (0,2) moduli, namely generically at least two complex Wilson line moduli  $B$  and  $C$  [22]. ( $B$  and  $C$  being, for example,  $\underline{27}$  respectively  $\underline{\bar{27}}$  of  $E_6$ ). Then the Kähler potential for these fields can be determined as follows [22]:

$$K_0 = -\log[(T_i + \bar{T}_i)(U_i + \bar{U}_i) - \frac{1}{2}(B_i + \bar{C}_i)(\bar{B}_i + C_i)]. \quad (3.13)$$

The corresponding Kähler moduli space is given by  $\mathcal{K}_0 = SO(2, 4)/SO(2) \times SO(4)$ . A few remarks are at hand. First note that in the absence of Wilson lines ( $B = C = 0$ ) the Kähler potential splits into the sum  $K_0 = K(T, \bar{T}) + K(U, \bar{U})$ , which is the well-known Kähler potential for the factorizable coset  $SO(2, 2)/SO(2) \times SO(2) = SU(1, 1)/U(1)_T \otimes SU(1, 1)/U(1)_U$ . On the other hand, turning on Wilson lines, the moduli space does not factorize anymore into two submanifolds. Thus it is natural to expect that also in a more general situation the moduli space is not anymore factorizable into a space of the Kähler class moduli times a space of the complex structure moduli (as it is true for (2,2) compactifications) as soon as (0,2) moduli are turned on. Also note that the complex Wilson lines give rise to holomorphic  $BC$  and antiholomorphic  $\bar{B}\bar{C}$  terms in the Kähler

potential. This is in principle just what is needed for the solution of the  $\mu$ -problem; upon identification of  $H_1$  with  $B$  and  $H_2$  with  $C$  the mass mixing term becomes [23, 22]

$$H_{12} = \frac{1}{(T + \bar{T})(U + \bar{U})}. \quad (3.14)$$

(This is also true in general if  $B$  and  $C$  are not moduli but matter fields with tree-level zero vev's [4, 23].) Thus we learn that holomorphic mixing terms in the Kähler potential can occur only if  $\theta_i = \pm 1$ , i.e. if there exists a complex structure modulus  $U_i$ . Consequently, the Higgs fields should be associated to this particular complex plane.

Now let us briefly discuss the duality symmetries. We consider the most interesting case with four complex moduli  $T, U, B$  and  $C$ . (For more discussion see [22].) In addition we assume that the complex plane corresponds to a two-dimensional subtorus. The duality, i.e. modular group in question is then given by the discrete group  $O(2, 4, Z)$ . The modular group  $O(2, 4, Z)$  contains an  $SO(2, 2, Z) = PSL(2, Z)_T \times PSL(2, Z)_U$  subgroup.  $PSL(2, Z)_T$  acts in the standard way on the  $T$  modulus

$$T \rightarrow \frac{aT - ib}{icT + d} \quad (3.15)$$

( $a, b, c, d \in Z, ad - bc = 1$ ). However  $U$  transforms also non-trivially under this transformation as

$$U \rightarrow U - \frac{ic}{2} \frac{BC}{icT + d}. \quad (3.16)$$

Thus, in the presence of  $B$  and  $C$ ,  $T$  and  $U$  get mixed under duality transformations [22, 23] which reflects the non-factorizable structure of the moduli space.

For the discussion of supersymmetry breaking one also needs to include one-loop corrections to the moduli Kähler potential. These arise due to a one-loop mixing of the  $S$ -field with the internal moduli. This is the so-called Green-Schwarz mixing with mixing coefficient  $\delta_{GS}^i$ . Specifically one can show that the loop corrected Kähler potential has the following structure [12, 24]:

$$K_0^{1-\text{loop}} = -\log Y + K_0^{\text{tree}}(T, U, A, B, C), \quad (3.17)$$

$$Y = S + \bar{S} + \frac{1}{8\pi^2} \sum_{i=1}^3 \delta_{GS}^i K_{0 \text{ tree}}(T_i, U_i, A_i, B_i, C_i).$$

Furthermore for the computation of the SSBP we need the tree-level Kähler potential of the matter fields. It can be shown to have the following form [25, 3]:

$$K_{\alpha\bar{\beta}} = \delta_{\alpha\bar{\beta}} \prod_{i=1}^3 (T_i + \bar{T}_i)^{n_\alpha^i}. \quad (3.18)$$

(For simplicity we have included only the generic  $T_i$  moduli.) The integers  $n_\alpha^i$  are called modular weights of the matter fields, since the  $Q_\alpha$  transform under  $PSL(2, Z)$  as

$$Q^\alpha \rightarrow Q^\alpha \prod_{i=1}^3 (ic_i T_i + d_i)^{n_\alpha^i}. \quad (3.19)$$



(The Wilson line moduli  $A, B, C$  have modular weight -1.)

As a final ingredient we have to specify the form of the gauge kinetic function in orbifold compactifications. Including one  $A$ -type modulus, the  $f$ -function in lowest order in  $A$  is given as

$$f(S, T_i, A)_a = S - \frac{1}{8\pi^2}(b_1 - b_0) \log[h(T_i)A] - \frac{1}{8\pi^2} \sum_{i=1}^3 (b'_a{}^i - \delta_{GS}^i) \log \eta(T_i)^2. \quad (3.20)$$

Here  $\eta(T_i)$  is the well-known Dedekind function and reflects the one-loop threshold contributions of momentum and winding states [26]. The  $A$  contribution corresponds to the mass thresholds [27, 28] of those fields  $Q^\alpha$  which get mass by a superpotential coupling to  $A$ :  $W \sim h(T_i)A Q^\alpha Q^\beta$ . If one assumes that all matter fields, that are charged under  $G_a$ , get a  $A$ -dependent masses one obtains  $b_0 = -3C(G_a)$ ,  $b_1 = -3C(G_a) + \sum_\alpha T_a(\alpha)$ . Then  $b'_a{}^i = -C(G_a) + \sum_\alpha T_a(\alpha)(1 + 2n_\alpha^i)$ . It is not difficult to verify the correct duality transformation behaviour of  $f$ .

Now let us apply these formulas to discuss some specific aspects of supersymmetry breaking in orbifold compactifications. Let us focus on the non-perturbative gaugino condensation in the hidden gauge sector  $a$ . The non-perturbative superpotential then reads

$$W_{\text{SUSY-breaking}} = \frac{e^{\frac{24\pi^2}{b_0}S} [h(T_i)A]^{3(b_0-b_1)/b_0}}{\prod_{i=1}^3 [\eta(T_i)]^{6(b'_a{}^i - \delta_{GS}^i)/b_0}}. \quad (3.21)$$

This leads to the following expression [28] for the scalar potential  $V$  using the one-loop corrected Kähler potential but neglecting for simplicity a possible  $A$  contribution, i.e.  $b_0 = b_1 = 3b'_a{}^i$  (the inclusion of  $A$  can be found in [29, 28]):

$$V = m_{3/2}^2 \left\{ \left| 1 - \frac{24\pi^2}{b_0} Y \right|^2 + \sum_{i=1}^3 \frac{Y}{8\pi^2 Y - \delta_{GS}^i} (1 - 3 \frac{\delta_{GS}^i}{b_0}) (T_i + \bar{T}_i)^2 |\hat{G}_2(T_i)|^2 - 3 \right\}. \quad (3.22)$$

The minimization of this scalar potential leads to the following results. First note that in case of complete Green-Schwarz cancellation, i.e.  $b_0 = 3\delta_{GS}^i$ , there is no  $T_i$  dependence in the potential (as well as in  $m_{3/2}$ ) and  $T_i$  still remains as a undetermined parameter. On the other hand, for  $3\delta_{GS}^i \neq b_0$ , the modulus  $T_i$  gets dynamically fixed. A specific analysis was performed in [15, 17] for the case  $\delta_{GS}^i = 0$  with the result that at the minimum  $T_i \sim 1.2$  supersymmetry gets spontaneously broken in the  $T_i$  sector since at that point  $G_{T_i} \neq 0$ . However there is an important caveat within this analysis since it used the assumption that at the minimum  $G_S = 0$ . In fact, the above potential, triggered by the gaugino condensate, has no stable minimum with respect to  $S$ . Therefore the dilaton dynamics has to be modified in order to justify this assumption. One way could be that there are gaugino condensates in more than one hidden gauge sector [30]. Then  $G_S = 0$  is rather generic, however several  $\beta$ -function coefficients have to be tuned in a careful way in order to get  $m_{3/2} \sim \mathcal{O}(1\text{TeV})$ . A different, very interesting possibility is that the non-perturbative dilaton dynamics is governed by the so-called  $S$ -duality [31, 32]. This means that the true non-perturbative string partition function is actually  $PSL(2, Z)$

invariant resp. covariant with respect to the  $S$ -field due to non-perturbative monopole-like configuration in target space. The simplest possibility within this context is that the partition function looks like [31]

$$Z \sim \frac{1}{(S + \bar{S})|\eta(S)|^4}. \quad (3.23)$$

In the effective field theory this could mean that the effective superpotential contains a term  $\eta(S)^{-2}$  instead of the ‘standard’  $e^S$  dependence. Such types of superpotentials possibly lead to  $G_S = 0$ . Finally one has to remark in this context that the cosmological constant tends to be non-vanishing within the non-perturbative scenario, which is very disturbing but probably reflects our ignorance about the exact supersymmetry breaking dynamics. (For a recent discussion about the cosmological constant see [6]; in [33] it has been argued that a negative cosmological constant after gaugino condensation might be a desirable feature, for the fully renormalized cosmological constant to vanish.)

Now, we could proceed to calculate the SSBP resulting from this type of superpotentials. For example the squark and slepton masses are obtained as a function of the modular weights  $n_\alpha$  [3]. At this stage it is very convenient to parametrize the unknown supersymmetry dynamics by some angle  $\tan \theta \sim \frac{G_S}{G_T}$  [5], i.e. the relative strength of the supersymmetry breaking in the  $S$  and  $T$  sectors. Then the exact form of the (perturbative or non-perturbative) superpotential is parametrized by  $\theta$  and  $m_{3/2}$ , and the form of the SSBP depends only on known perturbative quantities like  $K$ . Specifically the scalar masses have the form (assuming vanishing cosmological constant, the index  $i$  is suppressed now) [5]:

$$m_\alpha^2 = m_{3/2}^2 [1 + n_\alpha (1 - \frac{\delta_{GS}}{24\pi^2 Y})^{-1} \cos^2 \theta]. \quad (3.24)$$

For arbitrary values of  $\theta$  these SSBP are non-universal. However for  $\theta = \pi/2$ , i.e. the dilaton dominated supersymmetry breaking, the SSBP are in fact universal [34]. Finally, for the gaugino masses similar expressions can be derived. Concluding, it would be very interesting to test some of these features in future colliders.

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